

# Stationary Universe Model: Inputs and Outputs<sup>1</sup>

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This talk presents the recent progress achieved in collaboration with A.Linde and D.Linde<sup>1,2,3</sup> towards understanding the true nature of the global spatial structure of the universe as well as the most general stationary characteristics of its time-dependent state with eternally growing total volume.

In our opinion, the simplest and, simultaneously, the most general version of inflationary cosmology is the chaotic inflation scenario. It can be realized in all models where the other versions of inflationary theory can be realized (see [4] for detailed account). Several years ago it was realized that inflation in these theories has a very interesting property<sup>5,6</sup> which will be discussed in this talk. If the universe contains at least one inflationary domain of a size of horizon ( $h$ -region) with a sufficiently large and homogeneous scalar field  $\phi$ , then this domain will permanently produce new  $h$ -regions of a similar type. During this process the total physical volume of the inflationary universe (which is proportional to the total number of  $h$ -regions) will grow indefinitely.

Fortunately, some kind of stationarity may exist in many models of inflationary universe due to the process of the universe self-reproduction<sup>4</sup>. The properties of inflationary domains formed during the process of the self-reproduction of the universe do not depend on the moment of time at which each such domain is formed; they depend only on the value of the scalar fields inside each domain, on the average density of matter in this domain and on the physical length scale. This talk will describe what assumptions and approximations are made in order to get this picture of the universe and what consequences they bear.

Let us consider the simplest model of chaotic inflation based on the theory of a scalar

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field  $\phi$  minimally coupled to gravity, with the Lagrangian

$$L = \frac{1}{16\pi}R + \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - V(\phi) . \quad (1)$$

Here  $G = M_p^{-2} = 1$  is the gravitational constant,  $R$  is the curvature scalar,  $V(\phi)$  is the effective potential of the scalar field. The most important fact for inflationary scenario is that for most potentials  $V(\phi)$  (e.g., in all power-law  $V(\phi) = g_n\phi^n/n$  and exponential  $V(\phi) = ge^{\alpha\phi}$  potentials) there is an intermediate asymptotic regime of slow rolling of the field  $\phi$  and quasi-exponential expansion of the universe. This expansion (inflation) ends at  $\phi \sim \phi_e$  where the slow-rolling regime  $\ddot{\phi} \ll 3H(\phi)\dot{\phi}$  breaks down.

During the inflation all the inhomogeneities are stretched away and, if the evolution of the universe were governed solely by classical equations of motion, we would end up with extremely smooth geometry of the spatial section of the universe with no primordial fluctuations to initiate the growth of galaxies and large-scale structure. Fortunately, the same gravitational instability which causes the growth of galaxies during the Hot Big Bang era leads to the existence of the growing modes of vacuum fluctuations during the inflation. The wavelengths of all vacuum fluctuations of the scalar field  $\phi$  grow exponentially in the expanding universe. When the wavelength of any particular fluctuation becomes greater than  $H^{-1}$ , this fluctuation stops oscillating, and its amplitude freezes at some nonzero value  $\delta\phi(x)$  because of the large friction term  $3H\dot{\phi}$  in the equation of motion of the field  $\phi$ . The amplitude of this fluctuation then remains almost unchanged for a very long time, whereas its wavelength grows exponentially. Therefore, the appearance of such a frozen fluctuation is equivalent to the appearance of a classical field  $\delta\phi(x)$  that does not vanish after averaging over macroscopic intervals of space and time.

Because the vacuum contains fluctuations of all wavelengths, inflation leads to the creation of more and more perturbations of the classical field with wavelengths greater than  $H^{-1}$ . The average amplitude of such non-linearly modulated perturbations generated during an e-fold time interval  $H^{-1}$  is given by

$$|\delta\phi(x)| \approx \frac{H}{2\pi} . \quad (2)$$

If the field is effectively massless (i.e.  $m^2(\phi) \ll H^2(\phi)$ ), the amplitude of each frozen wave does not change in time. On the other hand, phases of each waves are random. Therefore, the sum of all waves at a given point fluctuates and experiences Brownian jumps in all directions in the space of fields.

The standard way of description of the stochastic behavior of the inflaton field during the slow-rolling stage is to coarse-grain it over  $h$ -regions and consider the effective equation of motion of the long-wavelength field<sup>7,8,9</sup>:

$$\frac{d}{dt}\phi = -\frac{V'(\phi)}{3H(\phi)} + \frac{H^{3/2}(\phi)}{2\pi}\xi(t) , \quad (3)$$

Here  $\xi(t)$  is the effective white noise generated by quantum fluctuations, which leads to the Brownian motion of the classical field  $\phi$ .

This Langevin equation corresponds to the following Fokker-Planck equation for the probability density  $P_c(\phi, t)$  to find the field  $\phi$  at a given point (which now means  $h$ -region) after time  $t$ :

$$\frac{\partial P_c}{\partial t} = \frac{\partial}{\partial \phi} \left( \frac{H^{3/2}(\phi)}{8\pi^2} \frac{\partial}{\partial \phi} (H^{3/2}(\phi) P_c) + \frac{V'(\phi)}{3H(\phi)} P_c \right), \quad (4)$$

The formal stationary solution ( $\partial P_c / \partial t = 0$ ) of equation (4) would be<sup>7</sup>

$$P_c \sim \exp \left( \frac{3}{8V(\phi)} \right), \quad (5)$$

Note, that the stationary solution (5) is equal to the square of the Hartle-Hawking wave function of the universe<sup>10</sup>. At first glance, this result is a direct confirmation of the Hartle-Hawking prescription for the wave function of the universe.

However, in all realistic cosmological theories, in which  $V(\phi) = 0$  at its minimum, the distribution (5) is not normalizable. The source of this difficulty can be easily understood: any stationary distribution may exist only due to a compensation of a classical flow of the field  $\phi$  downwards to the minimum of  $V(\phi)$  by the diffusion motion upwards. However, diffusion of the field  $\phi$  discussed above exists only during inflation, i.e. only for  $\phi \geq 1$ ,  $V(\phi) \geq V(1) \sim m^2 \sim 10^{-12}$  for  $m \sim 10^{-6}$ . Therefore (5) would correctly describe the stationary distribution  $P_c(\phi)$  in the inflationary universe only if  $V(\phi) \geq 10^{-12} \sim 10^{80}$  GeV in the absolute minimum of  $V(\phi)$ , which is, of course, absolutely unrealistic<sup>6</sup>.

It can be shown<sup>1,2,6</sup> that the solutions of this equation with the effect of “end of inflation” boundary properly taken into account are non-stationary (decaying). It is of no surprise because we didn’t take into account yet the complicated and inhomogeneous expansion of the whole universe with multiplying number of  $h$ -regions.

In order to describe the structure of the inflationary universe beyond one  $h$ -region (minisuperspace) approach one has to investigate the probability distribution  $P_p(\phi, t)$ , which takes into account the inhomogeneous exponential growth of the volume of domains filled by field  $\phi$ <sup>5,6</sup>. The close view on the process of the growth of the volume of inflationary universe reveals the strong resemblance with branching diffusion processes<sup>1,2,11</sup>, with the role of branching particles being played by  $h$ -regions and the the diffusing parameter being associated with the inflaton field inside an  $h$ -region.

A branching diffusion process is characterized by the independent diffusion of particles within the allowed region and their splitting at random times into several “daughter” particles which then continue to diffuse independently from their birthplace. The characteristic splitting time is related to the intensity of branching which, in addition to diffusion coefficient and drift term of ordinary diffusion, determines the time evolution of the distribution of particles in the space where they live.

To describe the distribution of the inflaton field in the whole universe and not only in one Hubble domain, one has to modify the Fokker-Planck equation by introducing the branching intensity term corresponding to creation of more and more new “particles” during the diffusion process<sup>1,2,12,13,14,15</sup>.

$$\frac{\partial P_p}{\partial t} = \frac{\partial}{\partial \phi} \left( \frac{H^{3/2}(\phi)}{8\pi^2} \frac{\partial}{\partial \phi} (H^{3/2}(\phi) P_p) + \frac{V'(\phi)}{3H(\phi)} P_p \right) + 3H(\phi) P_p \quad (6)$$

The mathematical model describing such behavior is the recently developed theory of a new type of branching diffusion processes<sup>11</sup> where “new” refers to diffusion taking place in the parameter space rather than the space where the particles live. Indeed, in our case the  $h$ -regions live in ordinary space, while the parameter associated with each  $h$ -region is the value of the coarse-grained inflaton field in it and this parameter lives on the space of homogeneous fields which is in the simplest case a segment of the real line  $[\phi_e, \phi_p]$ . Eq. (6) is, from this perspective, the forward Kolmogorov equation for the first moment of the generating functional of number of branching particles.

There are two main sets of questions which may be asked concerning such processes. First of all, one may be interested in the probability  $P_p(\phi, t)$  to find a given field  $\phi$  at a given time  $t$  under the condition that initial value of the field was equal to some  $\phi(t=0) = \phi_0$ . In what follows we will denote  $\phi_0$  as  $\chi$ . On the other hand, one may wish to know, what is the probability  $P_p(\chi, t)$  that the given final value of the field  $\phi$  (or the state with a given final density  $\rho$ ) appeared as a process of diffusion and branching of a domain containing some field  $\chi$ . Or, more generally, what is the typical history of a branching Brownian trajectory which ends up at a hypersurface of a given  $\phi$  (or a given  $\rho$ )?

It happens that if we do not take the Planck boundary seriously, this equation also doesn't have a true stationary solution (the solutions found by Nambu *et. al.*<sup>13</sup> are heavily concentrated at the super-Planckian densities which is just another way to state that there is no real stationarity). This conclusion, however, may fail if we will treat the Planck boundary more carefully. There are different reasons to do this:

1. Diffusion equations were derived in the semiclassical approximation for quantum scalar field which breaks down near the Planck energy density  $\rho_p \sim M_p^4 = 1$ .
2. Interpretation of the processes described by these equations is based on the notion of classical fields in a classical space-time, which is not applicable at densities larger than  $\rho_p \sim 1$  because of large fluctuations of metric at such densities. In particular, our interpretation of  $P_c$  and  $P_p$  as of probabilities to find classical field  $\phi$  in a given point at a given time does not make much sense at  $\rho > 1$ .

There is also another, rather intrinsic for inflationary model, reason to expect the existence of the stationary solutions: as we will argue now, inflation kills itself as the density approaches the Planck density  $\rho_p \sim M_p^4$ .

In our previous investigation we assumed that the vacuum energy density is given by  $V(\phi)$  and the energy-momentum tensor is given by  $V(\phi) g_{\mu\nu}$ . However, quantum fluctuations of the scalar field give the contribution to the average value of the energy momentum tensor, which does not depend on mass (for  $m^2 \ll H^2$ ) and is given by<sup>16</sup>

$$\langle T_{\mu\nu} \rangle = \frac{3 H^4}{32 \pi^2} g_{\mu\nu} = \frac{2}{3} V^2 g_{\mu\nu} . \quad (7)$$

The origin of this contribution is obvious. Quantum fluctuations of the scalar field  $\phi$  freeze out with the amplitude  $\frac{H}{2\pi}$  and the wavelength  $\sim H^{-1}$ . Thus, they lead to the gradient energy density  $(\partial_\mu \delta\phi)^2 \sim H^4$ .

An interesting property of eq. (7) is that the average value of the energy-momentum tensor of quantum fluctuations does not look like an energy-momentum tensor corresponding to the gradients of a sinusoidal wave. It looks rather like a renormalization of the vacuum energy-momentum tensor (it is proportional to  $g_{\mu\nu}$ ). This means, in particular, that after averaging over all possible outcomes of the process of generation of long-wave perturbations, the result (for  $\langle T_{\mu\nu} \rangle$ ) does not depend on the choice of the coordinate system. However, if we are interested in local events (i.e. we are averaging only over short wavelengths), we will see long-wave inhomogeneities produced by the “frozen” fluctuations of the scalar field. This does not lead to any interesting effects at  $V \ll 1$  ( $\phi \ll \phi_p$ ) since in this case  $V^2 \ll V$ .

However, at the density larger than the Planck density the situation becomes much more complicated. At  $V > 1$  the gradient energy density  $\sim V^2$  becomes larger than the potential energy density  $V(\phi)$ . A typical wavelength of perturbations giving the main contribution to the gradient energy is given by the size of the horizon,  $l \sim H^{-1}$ . This means that the inflationary universe at the Planck density becomes divided into many domains of the size of the horizon, density contrast between each of these domains being of the order of one. These domains evolve as separate mini-universes with the energy density dominated not by the potential energy density but by the energy density of gradients of the field  $\phi$ . Such domains drop out from the process of exponential expansion. Some of them may re-enter this process later, but many of them will collapse within the typical time  $H^{-1}$ .

This can be effectively described by imposing some kind of absorbing (or reflecting, or elastic screen type) boundary conditions on our process of branching diffusion at some  $\phi_b$  which should be close to the Planck boundary. Our calculations have shown<sup>1</sup> that if the boundary conditions do not permit the field  $\phi$  penetrate deeply into the domain  $\phi > \phi_p$ , the final results of our investigation do not depend on the type of the boundary conditions imposed (whether they are absorbing, reflecting, etc.). They depend only on the value of the field  $\phi_b$  where the boundary condition is to be imposed, and this dependence is rather trivial. We have argued that  $\phi_b \sim \phi_p$ . Therefore we will assume now that the function  $P_p(\phi, t|\chi)$  satisfies boundary conditions corresponding to disappearing particles at  $\phi_p$  treating  $\phi_p$  now just as a phenomenological parameter not necessarily corresponding to  $V(\phi_p) = 1$ .

The boundary conditions at the “end of inflation” boundary can be derived from the requirement of the continuity of the probability and its flux through that boundary<sup>1,2</sup>. Their

form suggests that the dependence of the physical probability distribution on the value of the initial field  $\chi$  near that boundary is close to square of the “tunneling wavefunction”<sup>17</sup>.

We have shown<sup>1,2</sup> that the asymptotic solution for  $P_p(\phi, t|\chi)$  (in the limit  $t \rightarrow \infty$ ) is given by

$$P_p(\phi, t|\chi) \rightarrow e^{\lambda_1 t} \psi_1(\chi) \pi_1(\phi) . \quad (8)$$

Here  $\psi_1(\chi)$  is the only positive eigenfunction of the equation:

$$\frac{1}{2} \frac{H^{3/2}(\chi)}{2\pi} \frac{\partial}{\partial \chi} \left( \frac{H^{3/2}(\chi)}{2\pi} \frac{\partial}{\partial \chi} \psi_s(\chi) \right) - \frac{V'(\chi)}{3H(\chi)} \frac{\partial}{\partial \chi} \psi_s(\chi) + 3H(\chi) \psi_s(\chi) = \lambda_s \psi_s(\chi) . \quad (9)$$

$\lambda_1$  is the corresponding (real) eigenvalue, and  $\pi_1(\phi)$  (invariant density of branching diffusion) is the eigenfunction of the conjugate operator with the same eigenvalue  $\lambda_1$ . We found that in realistic theories of inflation the typical time of relaxation to the asymptotic regime is extremely short. It is only about a few thousands Planck times, i.e. about  $10^{-40}$  sec. This means, that the normalized distribution

$$\tilde{P}_p(\phi, t|\chi) = e^{-\lambda_1 t} P_p(\phi, t|\chi) \quad (10)$$

rapidly converges to a time-independent function:

$$\tilde{P}_p(\phi|\chi) \equiv \tilde{P}_p(\phi, t \rightarrow \infty|\chi) = \psi_1(\chi) \pi_1(\phi) . \quad (11)$$

It is this stationary distribution that we were looking for. After some calculations we came to the following expression for  $P_p(\phi|\chi)$  (note that the functions  $\Phi(\cdot)$  are essentially the same, the only difference is the argument):

$$P_p(\phi|\chi) = \frac{C}{V^{3/4}(\phi)} \exp \left\{ \frac{3}{16 V(\phi)} - \frac{3}{16 V(\chi)} \right\} \Phi(z(\chi)) \Phi(z(\phi)) \quad (12)$$

Here  $C$  is the normalization constant. Using the WKB approximation, we have calculated<sup>1,2</sup>  $\Phi(z(\phi))$  for a wide class of potentials usually considered in the context of chaotic inflation, including potentials  $V \sim \phi^n$  and  $V \sim e^{\alpha\phi}$ .

The solution we have found features interesting properties. First of all, it is concentrated heavily at the highest allowed values of the inflaton field. Despite the promising exponential prefactor in (12) which looks like a combination of the Hartle-Hawking<sup>10</sup> and tunneling<sup>17</sup> wavefunctions, the dependence of  $\Phi(z(\phi))$  on  $\phi$  appeared to be even steeper than those exponents. This function falls exponentially (with the rate governed by the Planckian energies) towards the lower values of the inflaton field and strongly overwhelms the dependence of the exponent in front of it. Only near the “end of inflation” boundary the solution (12) reveals something familiar — the dependence on the initial field  $\chi$  becomes similar to the square of tunneling wavefunction (simultaneously, the corresponding dependence on the final field  $\phi$  cancels out). And, as we have already mentioned, the relaxation time is very small (which is also a consequence of the fact that the dynamics of the self-reproducing universe is governed

by the maximal possible energies). The stationary distribution found in<sup>1,2</sup> is not very sensitive to our assumptions concerning the concrete mechanism of suppression of production of inflationary domains with  $\phi \gtrsim \phi_p$ . We hope that these results may show us a way towards the complete quantum mechanical description of the stationary ground state of the universe.

At the first glance it seems that there can be no observable consequences of the stationarity of inflation and the picture of self-reproducing universe which we just described. Indeed, according to traditional inflationary paradigm as well as according to our understanding, all the visible universe is produced during the last stages of the inflaton field rollover, when its energy is already much below not only the Planck scale but also the scale  $V(\phi_*)$  where the quantum diffusion becomes more important than the classical slow rolling. Therefore it seems that the effects related to essentially Planckian energies should determine the structure of the universe at such exponentially large scales which will inevitably be far beyond the observational limits, thus leaving us with traditional description of the visible part of the universe.

There is, however, one important fact which follows from stationarity and which should alert us about the possibility that we are overlooking a non-trivial effect. The very essence of stationary picture implies that the rate of the volume growth is constant and governed by the highest accessible energies! However, if we look only at the local (produced no more than 60 e-folds prior to the end of inflation) part of the universe, we are forced to conclude that those e-folds were achieved at much slower rate corresponding to nearly end-of-inflation energies. The resolution of the controversy is in the fact that the high total rate of volume growth for most values of inflaton field is achieved not by direct local expansion but by expansion at the highest accessible energies and subsequent slow rollover.

In other words, one can say that once we started to talk about the volume weighted “physical” probability distribution  $P_p(\phi, t|\chi)$ , we should be consistent and consider other effects with same volume weighted measure. Then it should not be surprising that some events which are very improbable in usual measure become highly probable in this new measure.

For example, in the traditional inflationary paradigm one obtains the typical evolution of the inflaton field at the end of inflation to be given by classical slow rollover plus comparatively small random noise (2). This result is a consequence of considering only local evolution of the field, where any deviations from it are strongly suppressed in the measure corresponding to  $P_c(\phi, t)$  which can be written<sup>4</sup> as  $\exp\left(-\frac{2\pi^2(\delta\phi)^2}{H^2(\phi)}\right)$ . However, if we look at the problem from the volume weighted point of view, we discover that the typical trajectory of the inflaton is the one which maximizes the overall volume of the universe. In other words, although some trajectory may be suppressed by the local probability, if it gives large enough contribution to the growth of the volume of the universe, one will find many such trajectories in the unit of physical volume<sup>3</sup>.

Let us find the typical inflaton trajectory in volume weighted measure. If it was typical in this measure for inflaton to sit at highest energies up to the time when it has to start rolling

down in order to approach the end of inflation by the time of our observation, it might be even better for it to sit there for a little longer and then to roll down with slightly higher speed than it is suggested by the classical slow rollover, because this will enable the part of the universe under consideration to expand even longer with higher rate and to produce more volume. However, the only way to achieve rollover speeds greater than the classical is for the quantum fluctuations to add up with one sign to greater jumps than it is suggested by (2).

Let the extra time interval spent at highest energies be  $\tilde{\Delta}t$ . Then we win the volume by factor of  $e^{3H_{max}\tilde{\Delta}t}$ . However, to compensate for the lost time the inflaton has to jump at least once with the amplitude  $\tilde{\delta}\phi$  such that:

$$\tilde{\Delta}t(\phi) = \frac{\tilde{\delta}\phi}{\dot{\phi}} = \frac{n(\phi) \frac{H(\phi)}{2\pi}}{\dot{\phi}} = n(\phi) \Delta t(\phi) \quad (13)$$

where we introduced the factor  $n(\phi)$  by which the jump is amplified. Such jump is suppressed in probability by the factor  $e^{-\frac{1}{2}n^2(\phi)}$ . The jump at a given value of  $\phi$  occurs with such amplitude that maximizes the volume weighted probability:

$$\max \left[ \exp \left( 3H_{max}n(\phi)\Delta t - \frac{1}{2}n^2(\phi) \right) \right] \quad (14)$$

which gives the answer

$$n(\phi) = 3H_{max}\Delta t(\phi) \quad (15)$$

In fact, we have found<sup>3</sup> that the typical trajectory consists entirely of such subsequent jumps. We can relate the amplitude of the jump to the characteristics of the universe at a given length scale through the field-dependent value of the time  $\Delta t(\phi) = \delta\phi / \dot{\phi}$  of the conventional quantum jumps.

Recall the expressions for the amplitudes of scalar and tensor perturbations generated at given inflaton field  $\phi$  and therefore associated with a length scale through the usual slow rolling:

$$A_S^{pert}(\phi) = \left( \frac{\delta\rho}{\rho} \right)_S = c_S \frac{H(\phi)\delta\phi}{\dot{\phi}} \Big|_{k \sim H} \quad (16)$$

$$A_T^{pert}(\phi) = \left( \frac{\delta\rho}{\rho} \right)_T = c_T \frac{H(\phi)}{M_p} \Big|_{k \sim H} \quad (17)$$

from where one obtains:

$$\frac{A_S^{pert}(\phi)}{A_T^{pert}(\phi)} = \frac{c_S}{c_T} M_p \frac{\delta\phi}{\dot{\phi}} = \frac{c_S}{c_T} M_p \Delta t(\phi) \quad (18)$$

Using this result we can rewrite (15) in the form:

$$n(\phi) = 3 \frac{H_{max}}{M_p} \frac{c_T}{c_S} \frac{A_S^{pert}(\phi)}{A_T^{pert}(\phi)} \quad (19)$$



In the same way as the conventional amplitude of jumps  $H/2\pi$  is related with the well known perturbations of the background energy density the “nonperturbatively amplified” jumps which we have just described are related with “nonperturbative” contribution to deviations of the background energy density from its average value. The correct interpretation of this result is that at the length scale associated with the value of the field  $\phi$  there is an additional nonperturbative contribution to *monopole* amplitude:

$$A_S^{nonpert}(\phi) = \left( \frac{3c_T}{c_S} \frac{H_{max}}{M_p} \frac{A_S^{pert}(\phi)}{A_T^{pert}(\phi)} \right) A_S^{pert} \quad (20)$$

The reason why there is no contribution to higher multipoles is that large quantum jumps proceed through formation of spherically symmetric transition regions (all deviations from spherical symmetry are suppressed exponentially). However, one could justifiably worry about the induced dipole contribution which might come from the observation of a spherically symmetric distribution from a non-central point. We have found<sup>3</sup> reasonable constraints on such induced contribution.

The self-consistency constraints for these results are found<sup>3</sup> to be:

$$n(\phi) \gg 1 \iff \frac{A_S^{pert}}{A_T^{pert}} \gg 1 \quad (21)$$

$$n(\phi)H(\phi) \ll M_p \iff (\nabla\phi)^2 \ll V(\phi) \quad (22)$$

which do not lead to any new restrictions for all models which are otherwise good from the observational point of view. There is another constraint which is not mentioned here and which can sometimes become important, especially in the cases when  $\frac{A_S^{pert}}{A_T^{pert}}$  gets extremely large.

Since at any scale the nonperturbative contribution is spherically symmetric (monopole-like) and the sign of this contribution is definite (the jumps of inflaton field are definitely towards the lower values) one can argue that observationally such a configuration will look like a void with scale-dependent density which is lower than the density of the surroundings. Such a void will resemble in many aspects a locally open universe. Indeed, if there is a non-perturbative spherically symmetric hill in the Newtonian potential and we are near the top of that hill, we should see additional spherically symmetric radial component of peculiar velocities at which the galaxies are moving from us at farther distances as compared to the uniform Hubble law of flat universe — a picture not very different from the open universe, at least in what concerns the motion of the matter. However, it will be interesting to obtain other observational characteristics of such spatial distribution of Newtonian potential.

The amplification factor in (20) is small unless  $H_{max} \sim M_p$ , which tells us that one indeed needs to rely on the possibility of inflation going as high as Planckian energies in order to get such an effect. On the other hand it is proportional to the ratio of the scalar perturbative amplitude to tensor perturbative amplitude, which is a quite large number in all reasonable

inflationary theories, and may become very large in some versions of inflation with especially low end-of-inflation energy scale. For example, in the usual  $\lambda\phi^4$  chaotic inflation one gets  $\frac{A_S^{pert}}{A_T^{pert}} \sim 10$ ,  $H_{max} \sim M_p$ , and finally  $A_S^{nonpert} \sim 40 A_S^{pert} \sim 10^{-3}$ .

In some more exotic inflationary scenarios one can get a huge amplification factor  $\frac{A_S^{pert}}{A_T^{pert}}$  and our non-perturbative monopole contribution may actually become of order 1. This would be the best so far way to reconcile  $\Omega < 1$  with inflation retaining small anisotropy and other well established achievements of inflation. However, this is a case when one should be very careful about checking the compliance with self-consistency constraints which may invalidate the result. There are at least some scenarios which survive the first crude check and we are currently working on refining them.

As we promised in the title of this talk, we have outlined the inputs which go into building the stationary universe scenario, and quite surprisingly we have found some observational outputs which are so unique and unmistakable that they would be very difficult to mimic in any other model. They intrinsically depend on the assumption that the inflation can go up to Planckian energies and that it continues essentially forever. It becomes apparent that the picture of stationary universe is not only theoretically consistent and aesthetically appealing but is also testable and might be just the one needed to describe the observational data.

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